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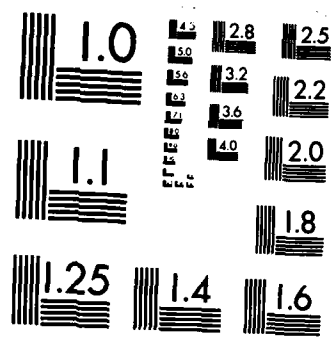
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Low Frequency Instabilities Driven by an Ion Temperature Anisotropy

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LOW FREQUENCY INSTABILITIES DRIVEN BY AN ION TEMPERATURE ANISOTROPY

I. INTRODUCTION

It is well known that HANE's can have deleterious effects on C^3I systems (for example, by disrupting satellite communications and radar systems) and it is therefore important to the Defense Nuclear Agency to understand the atmospheric effects of HANE's. The primary cause of these negative effects is the production of long-lasting, large-scale, ionization irregularities in the ionosphere. In order to understand the cause and evolution of these irregularities, the Naval Research Laboratory has developed and is continuing to develop detailed theoretical and computational models of HANE's. Of recent interest to DNA is a thorough re-examination of early time phenomena ($t \lesssim 1$ sec) which are primarily associated with the deposition of weapon energy in the atmosphere. Aside from radiation processes, the conversion of the kinetic energy of the debris into thermal energy of both the debris and air ions, and its ultimate deposition in the conjugate patches, is of particular importance. The important issues involved are debris-air coupling length and time scales, and thermalization processes. Underlying these issues are a host of complicated plasma physics processes. The ensuing turbulence which results from plasma instabilities can act as a means to thermalize the various ion species and to pitch angle scatter ions.

To date, the most intense effort to understand the role of plasma instabilities in early time HANE evolution has been directed at debris-air coupling. In this regard, NRL performed a series of investigations in the 70's to determine the anomalous transport properties of instabilities which could occur at the debris-air interface (Lampe et al., 1975). The effects of these instabilities on the evolution of the coupling shell were determined by incorporating the appropriate anomalous transport

Coefficients in a multi-fluid code (Clark, private communication). Subsequently, further advances have been made in this area through the use of a one dimensional hybrid code. This code treats the electrons as a fluid but the ions as particles. This permits the instabilities, and their effects on the plasma, to be studied self-consistently (Goodrich et al., 1984). It has also been recently suggested that microturbulence could impact the evolution of structure causing instabilities (Huba, 1984). Thus, the point to be made is that plasma turbulence can have a significant impact on the early time evolution of a HANE.

The aforementioned studies of plasma instabilities have dealt with high frequency turbulence, i.e., $\omega \gg \Omega_i$ where ω is the wave frequency and $\Omega_i = eB/m_i c$ is the ion cyclotron frequency. These instabilities impact very early time processes such as debris-air coupling and plasma thermalization. However, there are many low frequency instabilities (i.e., $\omega \lesssim \Omega_i$) that can be excited which may be important to early time HANE evolution. Specifically, the electromagnetic ion cyclotron instability ($\omega \sim \Omega_i$) has the capability of pitch angle scattering energetic ions (Davidson and Ogden, 1975). Debris and air ions are energized predominantly in a direction perpendicular to the ambient magnetic field B . Electromagnetic ion cyclotron waves, if excited, can then "transform" this perpendicular energy to parallel energy. This allows energetic ions to stream down the field lines into the upper atmosphere to deposit their energy; an important consideration in HANE events such as Starfish. Recent 1D particle simulations have observed this effect (Ambrosiano and Brecht, 1984). Aside from the electromagnetic ion cyclotron instability, other low frequency hydromagnetic instabilities may be excited which can be important. For example, the mirror or firehose

instabilities (Hasegawa, 1971) generate electromagnetic turbulence which can propagate away from the coupling shell into the magnetosphere. These waves are capable of pitch angle scattering ambient radiation belt protons which can lead to a redistribution of high energy protons in the magnetosphere (Cladis et al., 1970). This issue is germane to satellite survivability since it is believed that Starfish caused the enhancement of high energy proton fluxes at low L shells (Filz, 1967) which affected the operation of several satellites.

One mechanism to generate the instabilities described above is an anisotropic ion distribution. Specifically, a bi-Maxwellian type distribution such that $T_{\perp i} \neq T_{\parallel i}$ where $T_{\perp i}$ and $T_{\parallel i}$ are the perpendicular and parallel ion temperatures, respectively. The purpose of this report is to investigate in detail the linear stability properties of low frequency waves ($\omega \lesssim \Omega_i$) which are driven unstable by ion anisotropies. Specific application to HANE events will also be discussed.

The organization of the paper is as follows. In the next section we present the basic assumptions of the analysis and derive a general dispersion equation which describes low frequency modes ($\omega \lesssim \Omega_i$). In Section III we present both analytical and numerical results. Finally, in Section IV we summarize our results and apply them to HANE's.

II. DERIVATION OF DISPERSION EQUATION

We assume a spatially homogeneous plasma with an ion temperature anisotropy ($T_{\perp i} \neq T_{\parallel i}$) that is immersed in a uniform static magnetic field $\mathbf{B} = B_0 \hat{z}$. The perturbed electric field is taken to be $\mathbf{E}(\mathbf{x}, t) = \mathbf{E} \exp [i(\mathbf{k} \cdot \mathbf{x} - \omega t)]$ where $\mathbf{k} = k_{\perp} \hat{x} + k_{\parallel} \hat{z}$ without any loss of generality because of the symmetry of the unperturbed system. The perpendicular wavelength of

the perturbation is assumed to be much larger than the electron Larmor radius ρ_e ($k_\perp \rho_e \ll 1$) but can be comparable to the ion Larmor radius ρ_i ($k_\perp \rho_i \sim 1$). The frequency, as well as the growth rate, of the perturbation is taken to be of order of the ion cyclotron frequency ($\omega \lesssim \Omega_i$). For simplicity, we assume that the equilibrium ion distribution is described by a bi-Maxwellian distribution function

$$F_i(v_\perp^2, v_\parallel^2) = \left(\frac{m_i}{2\pi T_{i\perp}}\right)^{1/2} \left(\frac{m_i}{2\pi T_{i\parallel}}\right) \exp\left(-\frac{m_i v_\perp^2}{2T_{i\perp}} - \frac{m_i v_\parallel^2}{2T_{i\parallel}}\right), \quad (1)$$

where m_i is the mass of the ion species. The electron distribution function is assumed to be isotropic since collisions, as well as high frequency instabilities driven by any electron anisotropy, tend to remove the velocity-space anisotropy in the electron distribution on a very short time scale (i.e., $t \ll \Omega_i^{-1}$) [see Appendix].

Linearizing the Vlasov equation and using Maxwell's equation, a set of homogeneous equations $\mathbf{M} \cdot \mathbf{E} = 0$ describing the propagation of electromagnetic disturbances at an arbitrary angle with respect to the magnetic field can be obtained (Krall and Trivelpiece, 1973), where

$$\mathbf{M} = 1 - \frac{k^2 c^2 - c^2 \mathbf{k}_\perp \mathbf{k}_\perp}{\omega^2} + \sum_{\alpha=i,e} \frac{2\omega_p^2}{\omega \Omega_\alpha} \sum_n \langle \mathbf{p}_{\alpha n} \mathbf{Q}_{\alpha n} \rangle_\alpha \quad (2)$$

and

$$\langle \quad \rangle_\alpha \equiv 2\pi \int dv_\parallel \int 2v_\perp dv_\perp (\quad),$$

$$\mathbf{p}_{\alpha n} \equiv \frac{n\Omega_\alpha J_n}{k_\perp} \mathbf{A}_n \hat{x} + i v_\perp J_n' \mathbf{A}_n \hat{y} + v_\parallel J_n \mathbf{B}_n \hat{z},$$

$$Q_n \equiv \frac{n\Omega_\alpha J_n}{k_\perp} \hat{x} - i v_\perp J_n' \hat{y} + v_\parallel J_n \hat{z}$$

$$A_n \equiv \left[\left(1 - \frac{k_\parallel v_\parallel}{\omega} \right) \frac{\partial f_\alpha}{\partial v_\perp^2} + \frac{k_\parallel v_\parallel}{\omega} \frac{\partial f_\alpha}{\partial v_\parallel^2} \right] (\omega - n\Omega_\alpha - k_\parallel v_\parallel)$$

$$B_n \equiv A_n + (\partial f_\alpha / \partial v_\perp^2 - \partial f_\alpha / \partial v_\parallel^2) / \omega,$$

$J_n = J_n(k_\perp v_\perp / \Omega_\alpha)$ is the Bessel function of order n , $J_n'(x) = dJ_n/dx$, and $\Omega_\alpha = e_\alpha B_0 / m_\alpha c$ is the cyclotron frequency of species α (e : electrons and i : ions).

After performing the velocity space integration for both species and taking the appropriate small gyroradius limit for electrons, we obtain the dispersion equation

$$\det [\underline{M}] = 0. \quad (3)$$

The components of \underline{M} are

$$\begin{aligned} M_{xx} &= 1 - \frac{k_\parallel^2 c^2}{\omega^2} + \frac{1}{2} \frac{\omega_{pe}^2}{\omega^2} (G_{e1} + G_{e-1}) + \frac{\omega_{pi}^2}{\omega^2} \sum_n \frac{n^2 \Gamma_n}{x} G_n \\ M_{xy} &= -M_{yx} = -i \frac{\omega_{pe}^2}{2\omega^2} (G_{e1} - G_{e-1}) + i \frac{\omega_{pi}^2}{\omega^2} \sum_n n \Gamma_n' G_n \\ M_{xz} &= M_{zx} = \frac{k_\parallel k_\perp c^2}{\omega^2} + \frac{\omega_{pe}^2}{2\omega\Omega_e} (H_{e1} - H_{e-1}) + \frac{\omega_{pi}^2}{\omega\Omega_i} \frac{k_\perp}{k_\parallel} \sum_n \frac{n \Gamma_n}{x} H_n \\ M_{yy} &= 1 - \frac{k_\perp^2 c^2}{\omega^2} + \frac{\omega_{pe}^2}{2\omega^2} (G_{e1} + G_{e-1}) + \frac{\omega_{pi}^2}{\omega^2} \sum_n \left[(x \Gamma_n')' + \Gamma_n \right] G_n \\ M_{yz} &= -M_{zy} = -i \frac{\omega_{pe}^2}{\omega\Omega_e} \frac{k_\perp}{k_\parallel} H_{e0} + i \frac{\omega_{pi}^2}{\omega\Omega_i} \frac{k_\perp}{k_\parallel} \sum_n \Gamma_n' H_n \end{aligned} \quad (4)$$

$$M_{zz} = 1 - \frac{k_{\perp}^2 c^2}{\omega^2} + \frac{2\omega_{pe}^2}{k_{\parallel}^2 v_e^2} H_{e0} + \frac{\omega_{pe}^2}{\Omega_i^2} \frac{k_{\perp}^2}{k_{\parallel}^2} \sum_n \frac{\omega - n\Omega_i}{\omega} \frac{\Gamma_n}{x} H_n$$

where

$$G_{en} = \frac{\omega}{k_{\parallel} v_e} Z\left(\frac{\omega - n\Omega_e}{k_{\parallel} v_e}\right); H_{en} = 1 + \frac{\omega - n\Omega_e}{k_{\parallel} v_e} Z\left(\frac{\omega - n\Omega_e}{k_{\parallel} v_e}\right); x = \frac{k_{\perp}^2 v_{\perp i}^2}{2\Omega_i^2},$$

$$G_n = \left[1 - \frac{\omega - n\Omega_i}{\omega} \left(1 - \frac{T_{\perp i}}{T_{\parallel i}}\right)\right] \frac{\omega}{k_{\parallel} v_{\parallel i}} Z\left(\frac{\omega - n\Omega_i}{k_{\parallel} v_{\parallel i}}\right) - \left(1 - \frac{T_{\perp i}}{T_{\parallel i}}\right),$$

$$H_n = \left[\frac{T_{\perp i}}{T_{\parallel i}} + \frac{n\Omega_i}{\omega} \left(1 - \frac{T_{\perp i}}{T_{\parallel i}}\right)\right] \left[1 + \frac{\omega - n\Omega_i}{k_{\parallel} v_{\parallel i}} Z\left(\frac{\omega - n\Omega_i}{k_{\parallel} v_{\parallel i}}\right)\right]; \frac{T_{\perp i}}{T_{\parallel i}} = \frac{v_{\perp i}^2}{v_{\parallel i}^2},$$

v_e is the electron thermal velocity, $v_{\perp i}$ and $v_{\parallel i}$ are the ion perpendicular and parallel thermal velocity, respectively, $Z(\xi)$ is the plasma dispersion function, $\Gamma_n \equiv e^{-x} I_n(x)$, $I_n(x)$ is the modified Bessel function of order n , and $\Gamma'_n = d\Gamma_n/dx$.

III. ANALYSIS OF DISPERSION EQUATION

A. Electromagnetic Ion Cyclotron Instability

We first consider the excitation of ion cyclotron waves with $\omega \sim \Omega_i$.

For the case of parallel propagation ($k_{\perp} = 0$), the linear dispersion relation can be greatly simplified to give

$$\det [M] = D_+ D_- \left(1 + \frac{2\omega_{pe}^2}{k_{\parallel}^2 v_e^2} H_{e0} + \frac{2\omega_{pi}^2}{k_{\parallel}^2 v_i^2} H_0\right) = 0, \quad (5)$$

where

$$D_{\pm} = 1 - \frac{k_{\parallel}^2 c^2}{\omega^2} + \frac{\omega_{pi}^2}{\omega^2} G_{\pm 1} + \frac{\omega_{pe}^2}{\omega^2} G_{e\pm 1}.$$

Davidson and Ogden (1975) have shown that the transverse electromagnetic portion of (5) exhibits a strong instability in the presence of an ion temperature anisotropy such that $T_{\perp i} > T_{\parallel i}$. This instability has a characteristic frequency at maximum growth rate of $\omega_r \lesssim \Omega_i$ and $\gamma \sim (\beta_{\perp i}/2)^{1/2} \Omega_i$ where $\beta_{\perp i} = 8\pi m_i T_{\perp i}/B_0^2$, and the corresponding wavelength is such that $k_{\parallel} c/\omega_{pi} \gtrsim 1$. The electrostatic branch of the dispersion relation can be shown to be rigorously stable by the Nyquist technique.

The assumption that $k_{\perp} = 0$ corresponds to an "infinite" wavelength in the direction orthogonal to the ambient magnetic B_0 . For application to HANE's this assumption is violated because the coupling shell has finite dimensions perpendicular to B_0 . Thus, it is necessary to investigate the linear stability properties of ion cyclotron waves for oblique propagation using (3) and (4) if we are to apply this instability to HANE's. Since, in general, (3) and (4) are not amenable to analytical results we solve them numerically for a variety of parameters.

In Fig. 1 we show the growth rate γ/Ω_i vs. parallel wavenumber $ck_{\parallel}/\omega_{pi}$ for $\theta = 0^\circ, 20^\circ, 40^\circ, 60^\circ$, and 80° where $\theta = \tan^{-1}(k_{\perp}/k_{\parallel})$. The other parameters used are $\omega_{pi}/\Omega_i = 400$, $m_i/m_e = 1836$, $T_{\perp i}/T_{\parallel i} = 20$, $\beta_{\perp i} = 1.0$, and $T_e/T_{\perp i} = 0.1$. There are two major effects of oblique propagation on the excitation of electromagnetic cyclotron waves. First, the growth rate of the instability decreases as the waves become more oblique, i.e., as θ increases. However, even for θ as large as $\sim 50^\circ$ the growth rate is only reduced by a factor of two. Thus, the waves have substantial growth rates for a wide range of angles and are not confined to nearly parallel propagation. Second, in general, as θ increases the bandwidth of unstable modes in k_{\parallel} space decreases. The small k_{\parallel} cutoff remains roughly constant ($ck_{\parallel}/\omega_{pi} \sim 0.5$) for $0^\circ < \theta < 80^\circ$, but the large k_{\parallel} cutoff is reduced by

more than a factor of two. We note that the shape of the $\theta = 60^\circ$ growth rate is different from the others; namely, there is not a sharp falloff for modes with $k_{\parallel} > k_{\parallel m}$ where $k_{\parallel m}$ corresponds to the k_{\parallel} for maximum growth. There is almost a "plateau" in the growth rate in this region. This parameter regime corresponds to the Harris instability (Soper and Harris, 1965; Gary et al., 1976). The Harris instability only occurs for obliquely propagating waves and, in general, has a maximum growth rate less than the parallel propagating electromagnetic ion cyclotron instability (Gary et al., 1976).

In Fig. 2 we plot γ_m/Ω_i and $ck_{\parallel m}/\omega_{pi}$ vs. θ where γ_m denotes the maximum growth rate as a function of parallel wavenumber and $k_{\parallel m}$ is the corresponding wavenumber at maximum growth. The other parameters used are the same as in Fig. 1. These curves emphasize the major points indicated by Fig. 1: (1) maximum growth occurs for $\theta = 0^\circ$ (or $k_{\perp} = 0$, purely parallel propagating waves); (2) there is substantial wave growth for most angles of propagation; and (3) the wavelength of the fastest growing mode increases as the angle of propagation becomes more oblique (i.e., closer to perpendicular propagation).

B. MHD Instabilities

We now investigate low-frequency MHD instabilities i.e., $\omega \ll \Omega_i$. However, prior to presenting detailed numerical results, we first make several simplifying assumptions so that analytical results can be given. We assume that $\Omega_e, \Omega_i \gg \omega$, $k_{\parallel} v_e, k_{\parallel} v_i$ without any special ordering for $\omega/k_{\parallel} v_e$ and $\omega/k_{\parallel} v_i$. Making use of the approximation $\xi Z(\xi) \approx -1 - 1/2\xi^2$, for $\xi \gg 1$ the following limiting matrix elements can be derived:

$$\left(1 - \frac{\omega_s^2}{\omega^2}\right) \left[-\frac{\omega^2}{k_{\perp}^2 V_A^2} - \frac{\alpha}{\alpha} \frac{k_{\parallel}^2}{k^2} - \left(1 + \beta_{\perp} - \frac{T_{\perp e}}{T_{\parallel e}} \beta_{\perp e}\right) \frac{k_{\perp}^2}{k^2}\right] = \frac{\alpha_{\perp e}}{2} \left(\frac{k_{\perp}^2}{k^2}\right) \left(\frac{T_{\perp e}}{T_{\parallel e}}\right). \quad (A9)$$

Equation (A9) corrects the algebraic error made in Basu and Coppi (1982), as well as extending their result to include ion temperature anisotropy. The instability criterion in this case is

$$k_{\parallel}^2 (c_{s\parallel}^2 + \alpha V_A^2) + k_{\perp}^2 V_A^2 \left[1 + \beta_{\perp i} + \beta_{\perp e} \left(1 - \frac{T_{\perp e}}{2T_{\parallel e}}\right)\right] < 0. \quad (A10)$$

For $k_{\parallel}^2/k_{\perp}^2 \ll 2/\beta_{\parallel}$, one of the mode becomes unstable if

$$\frac{T_{\perp e}}{T_{\parallel e}} \beta_{\perp e} > 2(1 + \beta_{\perp}). \quad (A11)$$

The unstable mode is superficially similar to the wellknown MHD "mirror" instability, but since, as noted by Basu and Coppi (1982), "frozen-in condition" is not imposed, it is not identical to the mirror mode. The other mode is always stable, the reason that Basu and Coppi were able to obtain instability can be traced to the algebraic error they made.

In the other limit, $k_{\parallel} v_{\parallel i}, k_{\parallel} v_{\parallel e} \gg \omega$, we find

$$\bar{M}_{zy} = \frac{\omega_{pi}^2}{\omega \Omega_i} \frac{k_{\perp}}{k_{\parallel}} \left(\frac{T_{\perp e}}{T_{\parallel e}} - \frac{T_{\perp i}}{T_{\parallel i}}\right) \quad (A12a)$$

$$M_{yy} = D_{xx} - \frac{\alpha n_{\parallel}^2}{\alpha n_{\parallel}^2} - n_{\perp}^2 \left[1 - \beta_{\perp i} \left(\frac{T_{\perp i}}{T_{\parallel i}} - 1\right) - \beta_{\perp e} \left(\frac{T_{\perp e}}{T_{\parallel e}} - 1\right)\right] \quad (A12b)$$

$$D_{zz} = 1 + \frac{2\omega_{pe}^2}{k_{\parallel}^2 v_{\parallel e}^2} + \frac{2\omega_{pi}^2}{k_{\parallel}^2 v_{\parallel i}^2} \quad (A12c)$$

where $\bar{\alpha} = 1 - (\beta_{\parallel} - \beta_{\perp})/2$, $\beta_{\parallel} = \beta_{\parallel e} + \beta_{\parallel i}$, $\beta_{\perp} = \beta_{\perp e} + \beta_{\perp i}$, and $M_{xy} = M_{yx} = 0$ as before. Focusing our attention again to the limit $k_{\parallel} v_e \gg \omega \gg k_{\parallel} v_{\parallel i}$, we find

$$\bar{H}_{e0} = \frac{T_{\perp e}}{T_{\parallel e}} \left(1 + i\sqrt{\pi} \frac{\omega}{k_{\parallel} v_{\parallel e}} \right). \quad (A4)$$

Hence

$$\bar{M}_{zy} = \frac{\omega_{pi}^2}{\omega \Omega_i} \frac{k_{\perp}}{k_{\parallel}} \frac{T_{\perp e}}{T_{\parallel e}} \quad (A5)$$

$$M_{yy} = D_{xx} - \bar{\alpha} n_{\parallel}^2 - n_{\perp}^2 \left(1 + \beta_{\perp e} - \beta_{\perp e} \left(\frac{T_{\perp e}}{T_{\parallel e}} - 1 \right) \right) \quad (A6)$$

$$D_{zz} = 1 + \frac{2\omega_{pe}^2}{k_{\parallel}^2 v_{\parallel e}^2} - \frac{\omega_{pi}^2}{\omega^2}. \quad (A7)$$

Redefining $\omega_s = k_{\parallel} v_{\parallel e} (m_e/m_i)^{1/2} = k_{\parallel} c_{s\parallel}$ to be the parallel ion sound frequency, (8) can be simplified to

$$1 - \frac{\omega_s^2}{\omega^2} = \bar{\alpha} \frac{k_{\perp}^2 \rho_{s\parallel}^2}{\frac{\omega^2}{k_{\parallel}^2 v_A^2} - \bar{\alpha}} + \frac{\beta_{\perp e}}{2} \frac{(k_{\perp}^2/k^2)(T_{\perp e}/T_{\parallel e})}{\frac{\omega^2}{k_{\parallel}^2 v_A^2} - \bar{\alpha} \frac{k_{\parallel}^2}{k^2} - [1 + \beta_{\perp i} - \beta_{\perp e} \left(\frac{T_{\perp e}}{T_{\parallel e}} - 1 \right)] \frac{k_{\perp}^2}{k^2}}, \quad (A8)$$

where $\beta_{\perp e} = (2c_{s\parallel}^2/v_A^2)(T_{\perp e}/T_{\parallel e})$, and $\rho_{s\parallel}^2 = 2c_{s\parallel}^2/\Omega_i^2$.

Due to the smallness of $k_{\perp}^2 \rho_{s\parallel}^2$, (A8) can be simplified by ignoring the $k_{\perp}^2 \rho_{s\parallel}^2$ term, which results in the following coupled-mode equation

APPENDIX

To recover the results of Basu and Coppi (1982), we need to include the electron temperature anisotropy. This can be most simply done by changing H_{en} to

$$\bar{H}_{en} \equiv \left[\frac{T_{\perp e}}{T_{\parallel e}} + \frac{n\Omega_e}{\omega} \left(1 - \frac{T_{\perp e}}{T_{\parallel e}} \right) \right] \left[1 + \frac{\omega - n\Omega_e}{k_{\parallel} v_{\parallel e}} Z \left(\frac{\omega - n\Omega_e}{k_{\parallel} v_{\parallel e}} \right) \right] \quad (A1)$$

and using $\bar{G}_{en} = \omega(\omega - n\Omega_e)^{-1}(\bar{H}_{en} - 1)$ for G_{en} . In addition, the M_{yy} element of (4) should be modified as follows

$$M_{yy} = 1 - n^2 + \frac{\omega_{pe}^2}{2\omega^2} \left(\frac{2k_{\perp}^2 v_{\perp e}^2}{\Omega_e^2} \bar{G}_{e0} + \bar{G}_{e1} + \bar{G}_{e-1} \right) + \frac{\omega_{pi}^2}{\omega \Omega_i} \frac{k_{\perp}}{k_{\parallel}} \sum \frac{n \Gamma_n}{x} H_n. \quad (A2)$$

In the low frequency long wavelength limit, the matrix elements become

$$D_{xx} = M_{xx} + n_{\parallel}^2 \bar{\alpha} = 1 + \frac{\omega_{pi}^2}{\Omega_i^2} \quad (A3a)$$

$$M_{yy} = D_{xx} - \bar{\alpha} n_{\parallel}^2 - n_{\perp}^2 (1 - \beta_{\perp 1} G_0 - \beta_{\perp e} \bar{G}_{e0}) \quad (A3b)$$

$$\bar{M}_{zy} = \frac{\omega_{pi}^2}{\omega \Omega_i} \frac{k_{\perp}}{k_{\parallel}} (\bar{H}_{e0} - H_0) \quad (A3c)$$

$$M_{xz} = M_{zx} = n_{\perp} n_{\parallel} \bar{\alpha} \quad (A3d)$$

$$D_{zz} = M_{zz} + \bar{\alpha} n_{\parallel}^2 = 1 + \frac{2\omega_{pe}^2}{k_{\parallel}^2 v_{\parallel e}^2} \bar{H}_{e0} + \frac{2\omega_{pi}^2}{k_{\parallel}^2 v_{\parallel i}^2} H_0. \quad (A3e)$$

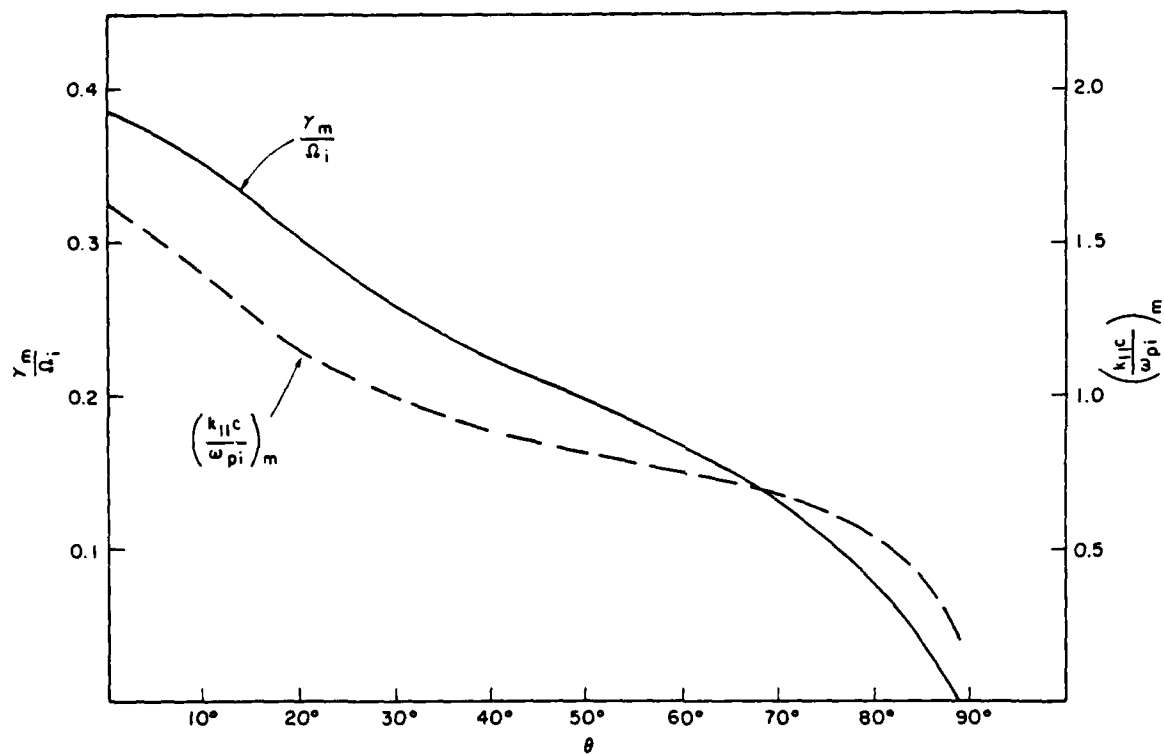


Fig. 2. Plot of γ_m/Ω_i and $ck_{||m}/\omega_{pi}$ versus θ where γ_m denotes the maximum growth rate as a function of parallel wavenumber and $k_{||m}$ is the corresponding wavenumber. The other parameters are the same as in Fig. 1.

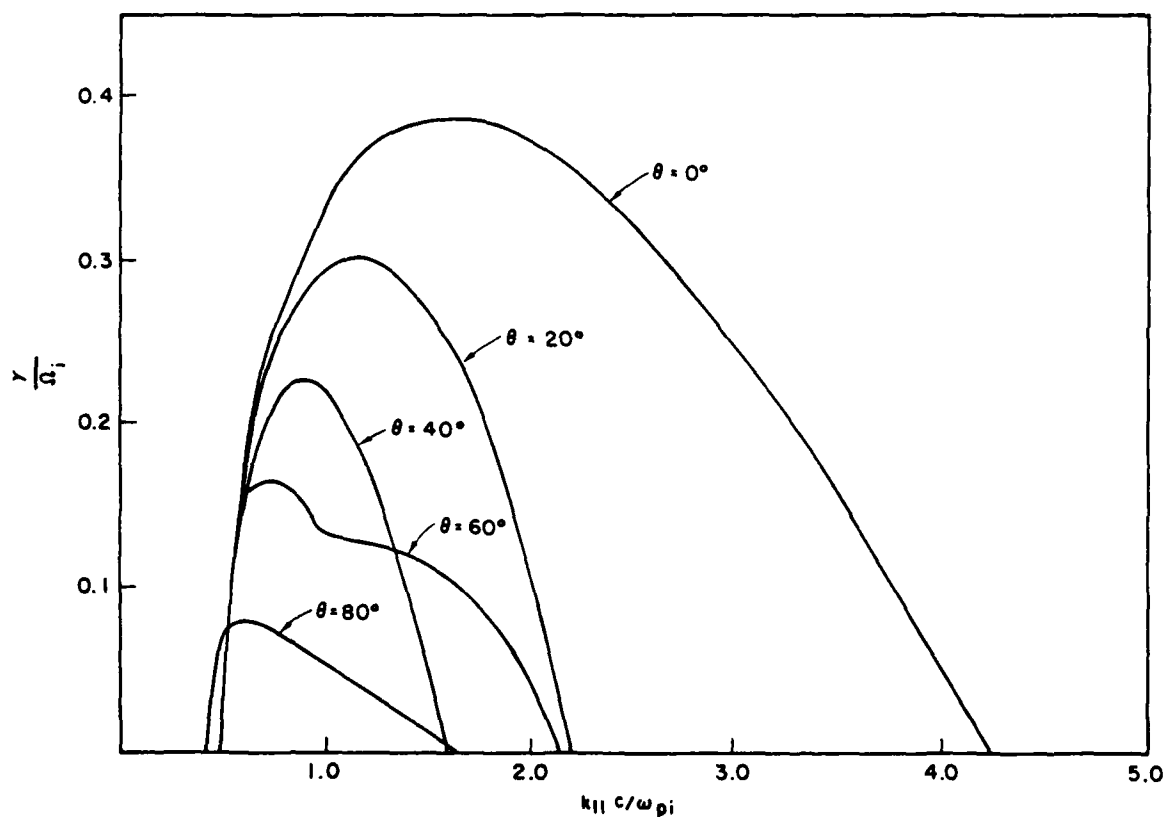


Fig. 1. Plot of γ/Ω_i versus $ck_{||}/\omega_{pi}$ for $\theta = 0^\circ, 20^\circ, 40^\circ, 60^\circ$, and 80° where $\theta = \tan^{-1}(k_{\perp}/k_{||})$. The other parameters used are $\omega_{pi}/\Omega_i = 400$, $m_i/m_e = 1836$, $T_{\perp i}/T_{\parallel i} = 20$, $v_{\perp i} = 1.0$, and $T_e/T_{\perp i} = 0.1$. Note that although the growth rate is a maximum for $\theta = 0^\circ$, it still has an appreciable value for oblique propagation.

is believed to have occurred following Starfish (Fitz, 1967). And second, these waves (i.e., electromagnetic ion cyclotron) have a significant component of δE_{\parallel} since they propagate obliquely, i.e., $k_{\perp} > k_{\parallel}$. These waves can be resonance with cold magnetospheric electrons which can be heated to 10-50 eV. This can lead to enhanced particle precipitation (Cornwall and Vesecky, 1984).

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applying these results to HANES a minimum time scale $\tau \sim \Omega_i^{-1}$ is required before the instabilities can be excited. However, one expects an instability to e fold several times before the fluctuating fields reach a sufficiently large amplitude to affect the particles. Thus, the time scale on which ion cyclotron waves become important is $\tau \sim 5\gamma^{-1} \sim 15 \Omega_i^{-1}$. Note that the ion cyclotron frequency is based on the value of the compressed magnetic field in the coupling shell. Since this value can be as much as an order of magnitude (or larger) than the ambient field, this time scale can be relatively fast. For example, assuming $m_i = 28$, $B_0 = 5G$, and $Z = 1$ we find that $\Omega_i^{-1} = 0.6$ msec so that $\tau \sim 9$ msec. Thus, in modeling the effects of electromagnetic ion cyclotron turbulence on debris and air ions (namely, pitch angle scattering) a minimum time scale is required which should be implemented in early time codes which model debris patch formation. We roughly estimate this time scale to be $\tau_m \sim 15\Omega_i^{-1}$ although it is parameter dependent and specific burst conditions need to be considered to accurately determine τ_m . We also add that pitch angle scattering is not the only mechanism which can convert perpendicular ion energy to parallel ion energy. The inverse mirror force associated with the compressed magnetic shell can also cause this energy conversion. The time scale with this process is shorter than pitch angle scattering since it is not associated with wave turbulence.

The waves discussed in this paper can have other important implications for HANE disturbances. First, these waves can also pitch angle scatter high energy ambient radiation belt protons ($E \gtrsim 1$ MeV). This can cause these high energy radiation belt particles to become trapped at lower L shells where defense and communication satellites reside. Enhanced fluences can then damage or incapacitate these satellites. In fact, this

anisotropy on the MHD instability and is an extension (and correction) of the work of Basu and Coppi (1982). Again, this is included for completeness. We now discuss the application of this research to HANE phenomena.

An important reason for this study was to investigate the stability properties of the electromagnetic ion cyclotron instability in order to determine whether or not it would be unstable in the HANE coupling shell. The significance of these unstable waves is that they pitch angle scatter ions, and therefore allow perpendicular energy to be transformed to parallel energy which is germane to debris patch formation. Most earlier studies focussed on parallel propagating modes which have "infinite" transverse wavelengths. However, the coupling shell has finite dimensions transverse to B_0 and one needs to consider oblique propagating modes for application to HANEs. The results of our analysis show that the electromagnetic ion cyclotron frequency has significant growth for a broad range of angles relative to B_0 . Although maximum growth occurs for $\theta = 0$ (i.e., $k_{\perp} = 0$), waves have relatively strong growth up to $\theta \sim 80^\circ$. This corresponds to transverse wavenumbers in the range $0 \lesssim k_{\perp} \lesssim 5 k_{\parallel}$; since $k_{\parallel} \sim \omega_{pi}/c$ (see Fig. 1) we can rewrite this inequality as $0 \lesssim k_{\perp} \lesssim 5\omega_{pi}/c$. Transforming to wavelength space we find that $\lambda_{\parallel} \sim c/\omega_{pi}$ and $\lambda_{\perp} \gtrsim c/5\omega_{pi}$. For HANE events such as Starfish and Checkmate, we note that the coupling shell width L is substantially larger than c/ω_{pi} at times $t \sim \text{few } \Omega_i^{-1}$, which implies $L \gg \lambda_{\perp}, \lambda_{\parallel}$. Thus, electromagnetic ion cyclotron waves will easily "fit" into the coupling shell, so that they can play a significant role in pitch angle scattering ions.

In this regard, another important issue is time scale. The instabilities considered in this report require magnetized ions; thus, in

through the small coupling coefficient ϵ . The preceding results derived for mirror instabilities correspond to weak coupling approximation of (10) by treating $D(\omega, k)$ to be of the order of ϵ . If we assume instead that $\omega^2 - \alpha k_{\parallel}^2 v_A^2 \sim O(\epsilon)$, then we arrive at the "firehose" dispersion relation (19). It is trivial to see that the firehose stability condition is $\beta_{\parallel i} - \beta_{\perp i} > 2$. In fact the stability condition remains unchanged as we change the ordering of ω , $k_{\parallel} v_e$, and $k_{\parallel} v_{\parallel i}$. Even with the inclusion of the electron temperature anisotropy, the stability condition merely changes to $\beta_{\parallel} - \beta_{\perp} > 2$. It therefore seems that the "firehose" mode is a very robust MHD mode which is insensitive to the kinetic effects associated with the parallel motion of electrons and ions. The firehose instability should not be a major concern for HANE since it required $T_{\parallel} > T_{\perp}$ which is not satisfied under ideal HANE conditions, at least in the coupling region during early time.

IV. DISCUSSION

We have presented a detailed linear analysis of low frequency ($\omega \lesssim \Omega_i$) instabilities driven by an ion temperature anisotropy ($T_{\perp i} \neq T_{\parallel i}$). In particular we have studied the electromagnetic ion cyclotron instability, the mirror instability, and the firehose instability. These instabilities require a high β plasma (i.e., $\beta \gtrsim 1$) which is typical of an early time HANE plasma. The electromagnetic ion cyclotron and mirror instabilities also require $T_{\perp i} > T_{\parallel i}$, while the firehose instability requires $T_{\parallel i} > T_{\perp i}$. In general, we expect the early time coupling shell to have $T_{\perp i} > T_{\parallel i}$ so that the firehose instability is not obviously relevant to HANE situations but has been included for the sake of completeness. We have also included an Appendix which describes the influence of an electron temperature

mode which are unstable under conditions discussed in their paper, are no longer so in our isotropic electron limit. In fact it is easy to see that (16) admits real ω^2 roots (i.e., they are either purely growing or pure oscillatory), and the sole condition for instability is

$$k_{\parallel}^2(c_s^2 + \alpha v_A^2) + k_{\perp}^2 v_A^2 (1 + \beta_{\perp i} + \beta_e/2) < 0, \quad (17)$$

which can only be satisfied for sufficiently negative values of α , or

$$\beta_{\parallel i} - \beta_{\perp i} > 2. \quad (18)$$

The instability condition (18) is the same as that of "firehose" instability even though the usual "frozen-in" condition is not imposed, as noted by Basu and Coppi (1982), and the polarization is different from that of the usual firehose mode. A more detailed discussion of this point is given in the Appendix.

To recover the "firehose" instability, we note that (10) can be rewritten in the following form

$$\left(\frac{\omega^2}{k_{\parallel}^2 v_A^2} - \alpha \right) D(\omega, \underline{k}) = \alpha k_{\perp}^2 \rho_s^2 \equiv \epsilon, \quad (19)$$

where $D(\omega, \underline{k}) = 0$ gives (11), the "mirror" mode dispersion equation. Equation (37) can be considered to be the linear mode-coupling equation between the mirror-type modes and the "firehose" mode

$$\omega^2 = \alpha k_{\parallel}^2 v_A^2 \equiv \left[1 - \frac{1}{2} (\beta_{\parallel i} - \beta_{\perp i}) \right] k_{\parallel}^2 v_A^2 \quad (20)$$

$$\bar{M}_{zy} = \frac{\omega_{pi}^2}{\omega \Omega_i} \frac{k_{\perp}}{k_{\parallel}} \quad (14a)$$

$$M_{zz} = 1 - n_{\perp}^2 \alpha + \frac{k_{De}^2}{k_{\parallel}^2} - \frac{\omega_{pi}^2}{\omega^2} \quad (14b)$$

$$M_{yy} = M_{xx} - n_{\perp}^2 (1 + \beta_{\perp i}), \quad (14c)$$

where we have made use of the identity $G_0 = H_0 - 1$ to approximate G_0 by -1 ,

$$M_{xx} = D_{xx} - \alpha n_{\parallel}^2, \quad D_{xx} = 1 + \omega_{pi}^2 / \Omega_i^2 \approx \omega_{pi}^2 / \Omega_i^2 \text{ and } M_{zz} = D_{zz} - \alpha n_{\perp}^2.$$

From (14b), $D_{zz} = 1 + k_{De}^2 / k_{\parallel}^2 - \omega_{pi}^2 / \omega^2 \approx (k_{De}^2 / k_{\parallel}^2) (1 - \omega_s^2 / \omega^2)$, where

$\omega_s = k_{\parallel} \omega_{pi} / k_{De} = k_{\parallel} c_s$ is the ion sound frequency. The dispersion relation (8) reduces to

$$1 - \frac{\omega_s^2}{\omega^2} = \alpha \frac{k_{\perp}^2 \rho_s^2}{\frac{\omega^2}{k_{\parallel}^2 v_A^2} - \alpha} + \frac{\beta_e}{2} \frac{k_{\perp}^2 / k^2}{\frac{\omega^2}{k_{\parallel}^2 v_A^2} - \frac{k_{\perp}^2}{k^2} (1 + \beta_{\perp i}) - \alpha \frac{k_{\parallel}^2}{k^2}} \quad (15)$$

Here again, the smallness of $k_{\perp}^2 \rho_s^2$ allows us to drop the first term on the right hand side of (15) provided that $|\omega^2 - \alpha k_{\parallel}^2 v_A^2| \gg k_{\perp}^2 \rho_s^2 \omega^2$. The resulting dispersion relation becomes

$$\left(1 - \frac{\omega_s^2}{\omega^2}\right) \left[\frac{\omega^2}{k_{\parallel}^2 v_A^2} - \alpha \frac{k_{\parallel}^2}{k^2} - (1 + \beta_{\perp i}) \frac{k_{\perp}^2}{k^2} \right] = \frac{\beta_e}{2} \frac{k_{\perp}^2}{k^2}. \quad (16)$$

Equation (16) closely resembles the equation derived by Basu and Coppi (1982) with one significant difference: their equation is derived under the assumption of large electron anisotropy and cold ions, whereas ours assumes isotropic electrons but retains temperature anisotropy for ions. Therefore the slow magnetosonic wave and the so-called "field-swelling"

where $\rho_s^2 = 2c_s^2/\Omega_i^2 = (v_e^2/\Omega_i^2)(m_e/m_i)$, $v_A^2 = c^2\Omega_i^2/\omega_{pi}^2$, $\beta_e = 2c_s^2/V_A^2 = 8\pi nT_e/B^2$ and $\beta_{ii} = 2V_{ii}^2/V_A^2 = 8\pi nT_{ii}/B^2$. For $T_e \sim T_{ii}$, T_{ii} , $\rho_s^2 \sim \rho_i^2$, we find $k_{\perp}^2 \rho_s^2 \ll 1$, consistent with our long wavelength ordering. The smallness of $k_{\perp}^2 \rho_s^2$ immediately yields, for frequencies significantly different from $\sqrt{\alpha} k_{\parallel} V_A$, the following approximate dispersion relation

$$(1 + \frac{T_e}{T_{ii}}) \{ \frac{\omega^2}{k^2 V_A^2} - \alpha \frac{k_{\parallel}^2}{k^2} - \frac{k_{\perp}^2}{k^2} [1 + \beta_{ii} (1 - \frac{T_{ii}}{T_{ii}})] \} = \frac{\beta_e}{2} (1 - \frac{T_{ii}}{T_{ii}})^2 \frac{k_{\perp}^2}{k^2}. \quad (11)$$

The stability condition is given by

$$\omega^2 = V_A^2 [1 - \frac{1}{2} (\beta_{ii} - \beta_{ii}) k_{\parallel}^2] + V_A^2 [1 + \beta_{ii} (1 - \frac{T_{ii}}{T_{ii}}) + \frac{\beta_e (1 - \frac{T_{ii}}{T_{ii}})^2}{2(1 + \frac{T_e}{T_{ii}})}] k_{\perp}^2 > 0 \quad (12)$$

For $k_{\parallel}^2/k^2 \ll 2/\beta_{ii}$, the stability condition becomes

$$1 + \beta_{ii} (1 - T_{ii}/T_{ii}) + \frac{T_e [T_{ii}^2 + T_{ii}^2] - 2 T_{ii} T_{ii}}{2 T_{ii} (T_{ii} + T_e)} > 0. \quad (13)$$

It is easy to see that instability can only occur when T_{ii} is sufficiently larger than T_{ii} . The stability condition (13) is the same as that of the MHD mirror mode (Hasegawa, 1971) Note that mirror instability only takes place when both k_{\parallel} and k_{\perp} are sufficiently nonzero, i.e., for oblique propagation.

2. Firehose Instability

In the limit $k_{\parallel} v_e \gg \omega \gg k_{\parallel} v_{ii}$, $H_{e0} = 1 + i/\pi\omega/k_{\parallel} v_e$, and $H_0 = -\frac{1}{2} k_{\parallel}^2 V_{ii}^2/\omega^2$. We arrive at the following approximate expressions for \bar{M}_{yz} , M_{zz} , and M_{yy} :

$$D_{zz} = \frac{\alpha n_{\perp}^2 D_{xx}}{D_{xx} - \alpha n_{\parallel}^2} + \frac{\bar{M}_{yz}^2}{M_{yy}}. \quad (8)$$

1. Mirror Instability

First, let us consider the familiar mirror instability. This corresponds to the limit $k_{\parallel} v_{\parallel i}, k_{\parallel} v_e \gg \omega$. In this limit D_{zz} becomes very large, hence $E_{\parallel}(E_z) \sim 0$ which is equivalent to "frozen-in condition" in ideal MHD treatment. Using small argument expansion for $Z(\xi_e)$ and $Z(\xi_i)$, we find

$$\bar{M}_{zy} = \frac{\omega_{pi}^2}{\omega \Omega_i} \frac{k_{\perp}}{k_{\parallel}} \left(1 - \frac{T_{ii}}{T_{\parallel i}}\right) \quad (9a)$$

$$M_{yy} = 1 + \frac{\omega_{pi}^2}{\Omega_i^2} - \alpha n_{\parallel}^2 - n_{\parallel}^2 \left[1 - \beta_{ii} \left(\frac{T_{ii}}{T_{\parallel i}} - 1\right)\right] \quad (9b)$$

$$D_{xx} = 1 + \frac{\omega_{pi}^2}{\Omega_i^2} \quad (9c)$$

$$D_{zz} = 1 + \frac{2\omega_{pe}^2}{k_{\parallel}^2 v_e^2} + \frac{2\omega_{pi}^2}{k_{\parallel}^2 v_{\parallel i}^2} \quad (9d)$$

Approximating $D_{zz} \approx (k_{De}^2/k_{\parallel}^2)(1 + T_e/T_{\parallel i})$, $k_{De}^2 = 2\omega_{pe}^2/v_e^2$, and $D_{xx} \approx \omega_{pi}^2/\Omega_i^2$, the dispersion relation (8) reduces to

$$1 + \frac{T_e}{T_{\parallel i}} = \alpha \frac{k_{\perp}^2 \rho_s^2}{\frac{\omega^2}{k_{\parallel}^2 v_A^2} - \alpha} + \frac{\beta_e}{2} \frac{(k_{\perp}^2/k_{\parallel}^2)(1 - T_{ii}/T_{\parallel i})^2}{\frac{\omega^2}{k_{\parallel}^2 v_A^2} - \alpha \frac{k_{\parallel}^2}{k^2} - \frac{k_{\perp}^2}{k^2} [1 + \beta_{ii}(1 - \frac{T_{ii}}{T_{\parallel i}})]}, \quad (10)$$

$$M_{xx} = 1 + \frac{\omega_{pi}^2}{\omega_i^2} - \frac{k_{\parallel}^2 c^2}{\omega^2} \left[1 - \frac{1}{2} (\beta_{\parallel i} - \beta_{\parallel e}) \right]$$

$$M_{xy} = M_{yx} = 0$$

$$M_{yy} = M_{xx} - \frac{k_{\perp}^2 c^2}{\omega^2} + \frac{k_{\perp}^2 c^2}{\omega^2} \beta_{\parallel i} G_0$$

(6)

$$M_{yz} = -M_{zy} = -i \frac{\omega_{pi}^2}{\omega \Omega_i} \frac{k_{\perp}}{k_{\parallel}} (H_{e0} - H_0)$$

$$M_{xz} = M_{zx} = \frac{k_{\perp} k_{\parallel}}{\omega^2} c^2 \left[1 - \frac{1}{2} (\beta_{\parallel i} - \beta_{\parallel e}) \right]$$

$$M_{zz} = 1 - \frac{k_{\perp}^2 c^2}{\omega^2} \left[1 - \frac{1}{2} (\beta_{\parallel i} - \beta_{\parallel e}) \right] + \frac{2\omega_{pe}^2}{k_{\parallel}^2 v_e^2} H_{e0} + \frac{2\omega_{pi}^2}{k_{\parallel}^2 v_i^2} H_0$$

where $\beta_{\parallel i} = \omega_{pi}^2 v_{\parallel i}^2 / \Omega_i^2 c^2$, $\beta_{\parallel e} = \omega_{pe}^2 v_{\parallel e}^2 / \Omega_e^2 c^2 = 8\pi n T_{\parallel i} / B^2$. In deriving (6) we have omitted terms which are smaller by order ω/Ω_i and ω/Ω_e . In addition, we have taken the small ion Larmor radius limit ($k_{\perp} \rho_{\perp i} \ll 1$). For $\beta_e, \beta_i \approx 1$ we can also neglect the 1 in M_{xx} as well as M_{zz} since $\omega_{pi}^2/\Omega_i^2, k_{\parallel}^2 c^2/\omega^2, k_{\perp}^2 c^2/\omega^2 \gg 1$. We now show how the well known mirror and firehose instabilities are recovered when subsidiary orderings are made.

Writing $M_{zy} = i\bar{M}_{zy}$ to make explicit the i dependence, we can express the dispersion relation (6) in the following compact form:

$$M_{zz} = \frac{n_{\perp}^2 n_{\parallel}^2 \left[1 - \frac{1}{2} (\beta_{\parallel i} - \beta_{\parallel e}) \right]^2}{M_{xx}} + \frac{\bar{M}_{yy}^2}{M_{yy}}, \quad (7)$$

where $n_{\perp}^2 \equiv c^2 k_{\perp}^2 / \omega^2$ and $n_{\parallel}^2 \equiv c^2 k_{\parallel}^2 / \omega^2$. Further simplification is possible with the introduction of $\alpha \equiv 1 - \frac{1}{2} (\beta_{\parallel i} - \beta_{\parallel e})$, $D_{xx} \equiv M_{xx} + \alpha n_{\parallel}^2$, and $D_{zz} \equiv M_{zz} + \alpha n_{\perp}^2$.

The "firehose" branch in this case still remains the same, but the coupled-mode equation becomes

$$\left[\frac{\omega^2}{k^2 v_A^2} - \frac{\alpha}{k^2} \frac{k_{\parallel}^2}{k^2} - \left(1 + \beta_{\perp} - \beta_{\perp i} \frac{T_{\perp i}}{T_{\parallel i}} - \beta_{\perp e} \frac{T_{\perp e}}{T_{\parallel e}} \right) \frac{k_{\parallel}^2}{k^2} \right] = \frac{\beta_{\parallel e} \left(\frac{k_{\perp}^2}{k^2} \right)}{2 \left(1 + \frac{T_{\parallel e}}{T_{\parallel i}} \right)} \left(\frac{T_{\perp e}}{T_{\parallel e}} - \frac{T_{\perp i}}{T_{\parallel i}} \right)^2. \quad (\text{A13})$$

Notice that the parallel ion-sound branch is gone because $k_{\parallel} v_{\parallel i} \gg \omega$, which means that both electrons and ions respond adiabatically to parallel electric field. The stability condition is given by

$$\omega^2 = v_A^2 \left[1 - \frac{1}{2} (\beta_{\parallel} - \beta_{\perp}) \right] k_{\parallel}^2 + v_A^2 \left[1 + \beta_{\perp} - \beta_{\perp i} \frac{T_{\perp i}}{T_{\parallel i}} - \beta_{\perp e} \frac{T_{\perp e}}{T_{\parallel e}} + \frac{\beta_{\parallel e} \left(\frac{T_{\perp e}}{T_{\parallel e}} - \frac{T_{\perp i}}{T_{\parallel i}} \right)}{2 \left(1 + \frac{T_{\parallel e}}{T_{\parallel i}} \right)} \right] k_{\perp}^2 > 0. \quad (\text{A14})$$

For $k_{\parallel}^2/k^2 \ll 2/\beta_{\parallel}$, the instability condition is the same as that of the MHD mirror instability (Hasegawa, 1971). Note that "frozen-in" condition is satisfied when $k_{\parallel} v_{\parallel i}, k_{\parallel} v_{\parallel e} \gg \omega$ since the perturbed parallel electric field is completely shielded out by the parallel electron and ion adiabatic response.

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